A comparative study of dark matter in the MSSM and its singlet extensions: a mini review

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Abstract

In this note we briefly review the recent studies of dark matter in the MSSM and its singlet extensions: the NMSSM, the nMSSM, and the general singlet extension. Under the new detection results of CDMS II, XENON, CoGeNT and PAMELA, we find that (i) the latest detection results can exclude a large part of the parameter space which allowed by current collider constraints in these models. The future SuperCDMS and XENON can cover most of the allowed parameter space; (ii) the singlet sector will decouple from the MSSM-like sector in the NMSSM, however, singlet sector makes the nMSSM quite different from the MSSM; (iii) the NMSSM can allow light dark matter at several GeV exists. Light CP-even or CP-odd Higgs boson must be present so as to satisfy the measured dark matter relic density. In case of the presence of a light CP-even Higgs boson, the light neutralino dark matter can explain the CoGeNT and DAMA/LIBRA results; (iv) the general singlet extension of the MSSM gives a perfect explanation for both the relic density and the PAMELA result through the Sommerfeld-enhanced annihilation. Higgs decays in different scenario are also studied.

PACS numbers: 14.80.Ly,11.30.Pb,95.35.+d

I. INTRODUCTION

Although there are many theoretical or aesthetical arguments for the necessity of TeV-scale new physics, the most convincing evidence is from the WMAP (Wilkinson Microwave Anisotropy Probe) observation of the cosmic cold dark matter, which naturally indicates the existence of WIMPs (Weakly Interacting Massive Particle) beyond the prediction of the Standard Model (SM). By contrast, the neutrino oscillations may rather imply trivial new physics (plainly adding right-handed neutrinos to the SM) or new physics at some very high see-saw scale unaccessible to any foreseeable colliders. Therefore, the TeV-scale new physics to be unraveled at the Large Hadron Collider (LHC) is the most likely related to the WIMP dark matter.

If WIMP dark matter is chosen by nature, it will give a strong argument for low-energy supersymmetry (SUSY) with R-parity which can give a good candidate. Nevertheless, SUSY is motivated for solving the hierarchy problem elegantly. It can also solve other puzzles of the SM, such as the 3σ deviation of the muon anomalous magnetic moment from the SM prediction. In the framework of SUSY, the most intensively studied model is the minimal supersymmetric standard model (MSSM) [1], which is the most economical realization of SUSY. However, this model suffers from the μ-problem. The μ-parameter is the only dimensional parameter in the SUSY conserving sector. From a top down view, one would expect the μ to be either zero or at the Planck scale. But in the MSSM, the relation of the electro-weak (EW) scale soft parameters ($\tilde{m}_d^2,\ \tilde{m}_u^2$) [2]

$$\frac{1}{2} M_Z^2 = \frac{\tilde{m}_d^2 - \tilde{m}_u^2 \tan^2\beta}{\tan^2\beta - 1} - \mu^2, \tag{1}$$

makes that μ must be at the EW scale, while LEP constraints on the chargino mass require μ to be non-zero [3]. A simple solution is to promote μ to a dynamical field in extensions of the MSSM that contain an additional singlet superfield $\hat{S}$ which does not interact with the MSSM fields other than the two Higgs doublets. An effective μ can be reasonably got at EW scale when $\hat{S}$ denotes the vacuum expectation value (VEV) of the singlet field. Among these extension models the next-to-minimal supersymmetric model (NMSSM) [4] and the nearly minimal supersymmetric model (nMSSM) [5, 6] caused much attention recently. Note that the little hierarchy problem which is also a trouble of the MSSM is relaxed greatly in the NMSSM.

If introduce a singlet superfield to the MSSM, the Higgs sector will have one more CP even component and one more CP odd component, and the neutralino sector will have one more singlino component. These singlet multiplets compose a "singlet sector" of the MSSM. It can make the phenomenologies of SUSY dark matter and Higgs different from the MSSM. More and more precision results of dark matter detection give us an opportunity to test if this singlet sector really exists. For example, experiments for the underground direct detection of cold dark matter $\tilde{\chi}$ have recently made significant progress. While the null observation of $\tilde{\chi}$ in the CDMS and XENON100 experiments has set rather tight upper limits on the spin-independent (SI) cross section of $\tilde{\chi}$-nucleon scattering [7, 8]. The CoGeNT experiment [9] reported an excess which cannot be explained by any known background sources but seems to be consistent with the signal of a light $\tilde{\chi}$ with mass around 10 GeV and scattering rate $(1\text{–}2) \times 10^{-40}$ cm^2. Intriguingly, this range of mass and scattering rate are compatible with dark matter explanation for both the DAMA/LIBRA data and the preliminary CRESST data [10]. Though CoGeNT result is not consistent with the CDMS or XENON results, it implies that the mass of dark matter can range a very long scope at EW scale, that is from a few GeV to several TeV. The indirect detection PAMELA also observed an excess of the cosmic ray positron in the energy range 10-100 GeV [11] which may be explained by dark matter.

In this paper, We will give a short review on the difference between the MSSM and the MSSM with a singlet sector under the constraints of new dark matter detection results. As the Higgs hunting on colliders has delicate relation with dark matter detections, the implication on Higgs searching is also reviewed. The content is based on our previous work [12–14]. the paper is organized as following, in sec. II, we will give a short review on the structures of the MSSM, the NMSSM and the nMSSM. In sec. III we will give a comparison on the models under the constraints of CDMS, XENON, and CoGeNT. In sec. IV, a general singlet extension of the MSSM is discussed, and a summary is given in sec. V.

II. THE MSSM AND ITS SINGLET EXTENSIONS

As the economical realization of supersymmetry, the MSSM has the minimal content of particles, while the NMSSM and the nMSSM extend the MSSM by only adding one singlet

Higgs superfield $\hat{S}$. The difference between these models is reflected in their superpotential:

$$\begin{aligned} W_{\rm MSSM} &= W_F + \mu \hat{H}_u \cdot \hat{H}_d, & (2)\\ W_{\rm NMSSM} &= W_F + \lambda \hat{H}_u \cdot \hat{H}_d \hat{S} + \frac{1}{3}\kappa \hat{S}^3, & (3)\\ W_{\rm nMSSM} &= W_F + \lambda \hat{H}_u \cdot \hat{H}_d \hat{S} + \xi_F M_n^2 \hat{S}, & (4) \end{aligned}$$

where $W_F = Y_u \hat{Q} \cdot \hat{H}_u \hat{U} - Y_d \hat{Q} \cdot \hat{H}_d \hat{D} - Y_e \hat{L} \cdot \hat{H}_d \hat{E}$ with $\hat{Q}$, $\hat{U}$ and $\hat{D}$ being the squark superfields, and $\hat{L}$ and $\hat{E}$ being the slepton superfields. $\hat{H}_u$ and $\hat{H}_d$ are the Higgs doublet superfields, λ, κ and ξ_F are dimensionless coefficients, and μ and M_n are parameters with mass dimension. Note that there is no explicit μ-term in the NMSSM or the nMSSM, and an effective μ-parameter (denoted as $\mu_{\rm eff}$) can be generated when the scalar component (S) of $\hat{S}$ develops a VEV. Also note that the nMSSM differs from the NMSSM in the last term with the trilinear singlet term $\kappa \hat{S}^3$ of the NMSSM replaced by the tadpole term $\xi_F M_n^2 \hat{S}$. As pointed out in Ref. [5], such a tadpole term can be generated at a high loop level and naturally be of the SUSY breaking scale. The advantage of such replacement is the nMSSM has no discrete symmetry thus free of the domain wall problem which the NMSSM suffers from.

Corresponding to the superpotential, the Higgs soft terms in the scalar potentials are also different between the three models (the soft terms for gauginos and sfermions are the same thus not listed here)

$$\begin{aligned} V_{\rm soft}^{\rm MSSM} &= \tilde{m}_d^2 |H_d|^2 + \tilde{m}_u^2 |H_u|^2 + (B\mu H_u \cdot H_d + {\rm h.c.}) & (5)\\ V_{\rm soft}^{\rm NMSSM} &= \tilde{m}_d^2 |H_d|^2 + \tilde{m}_u^2 |H_u|^2 + \tilde{m}_s^2 |S|^2 + \left(A_\lambda \lambda S H_d \cdot H_u + \frac{\kappa}{3} A_\kappa S^3 + {\rm h.c.} \right), & (6)\\ V_{\rm soft}^{\rm nMSSM} &= \tilde{m}_d^2 |H_d|^2 + \tilde{m}_u^2 |H_u|^2 + \tilde{m}_s^2 |S|^2 + \left(A_\lambda \lambda S H_d \cdot H_u + \xi_S M_n^3 S + {\rm h.c.} \right). & (7) \end{aligned}$$

After the scalar fields H_u,H_d and S develop their VEVs v_u, v_d and s respectively, they can be expanded as

$$H_d = \begin{pmatrix} \frac{1}{\sqrt{2}} (v_d + \phi_d + i\varphi_d) \\ H_d^- \end{pmatrix}, H_u = \begin{pmatrix} H_u^+ \\ \frac{1}{\sqrt{2}} (v_u + \phi_u + i\varphi_u) \end{pmatrix}, S = \frac{1}{\sqrt{2}} (s + \sigma + i\xi) . \quad (8)$$

The mass eigenstates can be obtained by unitary rotations

$$\begin{pmatrix} h_1 \\ h_2 \\ h_3 \end{pmatrix} = U^H \begin{pmatrix} \phi_d \\ \phi_u \\ \sigma \end{pmatrix}, \quad \begin{pmatrix} a_1 \\ a_2 \\ G_0 \end{pmatrix} = U^A \begin{pmatrix} \varphi_d \\ \varphi_u \\ \xi \end{pmatrix}, \quad \begin{pmatrix} G^+ \\ H^+ \end{pmatrix} = U^{H^+} \begin{pmatrix} H_d^+ \\ H_u^+ \end{pmatrix}, \quad (9)$$

where $h_{1,2,3}$ and $a_{1,2}$ are respectively the CP-even and CP-odd neutral Higgs bosons, G^0 and G^+ are Goldstone bosons, and H^+ is the charged Higgs boson. Including the scalar part of the singlet sector in the NMSSM and the nMSSM leads to a pair of charged Higgs bosons, three CP-even and two CP-odd neutral Higgs bosons. In the MSSM, we only have two CP-even and one CP-odd neutral Higgs bosons in addition to a pair of charged Higgs bosons.

The MSSM predicts four neutralinos χ_i^0 $(i = 1, 2, 3, 4)$, i.e. the mixture of neutral gauginos (bino λ' and neutral wino λ^3) and neutral higgsinos $(\psi^0_{H_u}, \psi^0_{H_d})$, while the NMSSM and the nMSSM predict one more neutralino corresponding to the singlino ψ_S from the fermion part of singlet sector. In the basis $(-i\lambda', -i\lambda^3 2, \psi^0_{H_u}, \psi^0_{H_d}, \psi_S)$ (for the MSSM ψ_S is absent) the neutralino mass matrix is given by

$$\begin{pmatrix} M_1 & 0 & m_Z s_W s_b & -m_Z s_W c_b & 0 \\ 0 & M_2 & -m_Z c_W s_b & m_Z c_W c_b & 0 \\ m_Z s_W s_b & -m_Z s_W s_b & 0 & -\mu & -\lambda v c_b \\ -m_Z s_W c_b & -m_Z c_W c_b & -\mu & 0 & -\lambda v s_b \\ 0 & 0 & -\lambda v c_b & -\lambda v s_b & \begin{cases} 2\frac{\kappa}{\lambda}\mu & \text{for the NMSSM} \\ 0 & \text{for the nMSSM} \end{cases} \end{pmatrix}, \tag{10}$$

where M_1 and M_2 are respectively $U(1)$ and $SU(2)$ soft gaugino mass parameters, $s_W = \sin\theta_W$, $c_W = \cos\theta_W$, $s_b = \sin\beta$ and $c_b = \cos\beta$ with $\tan\beta \equiv v_u/v_d$. The lightest neutralino $\tilde{\chi}^0_1$ is assumed to be the lightest supersymmetric particle (LSP), serving as the SUSY dark matter particle. It is composed by

$$\tilde{\chi}^0_1 = N_{11}(-i\lambda') + N_{12}(-i\lambda^3) + N_{13}\psi^0_{H_u} + N_{14}\psi^0_{H_d} + N_{15}\psi_S, \tag{11}$$

where N is the unitary matrix (N_{15} is zero for the MSSM) to diagonalize the mass matrix in Eq. (10). For the mass matrices above we should note that the following two points

1. For a moderate value of κ, the neutralino sector of the NMSSM can go back to the MSSM when λ approaches to zero. This is because in such case the singlino component will become super heavy and decouple from EW scale. The singlet scalar will not mix with the two Higgs doublet, then the NMSSM will be almost the same as the MSSM at EW scale.

2. Since the $\psi_S\psi_S$ element of Eq. (10) is zero in the nMSSM, the singlino will not decouple when λ approaches to zero. In fact, in the nMSSM the mass of the LSP can be written as

$$m_{\chi_1^0} \simeq \frac{2\mu_{\rm eff}\lambda^2(v_u^2+v_d^2)}{2\mu_{\rm eff}^2+\lambda^2(v_u^2+v_d^2)}\frac{\tan\beta}{\tan^2\beta+1}. \tag{12}$$

This formula shows that to get a heavy $\tilde{\chi}_1^0$, we need a large λ, a small $\tan\beta$ as well as a moderate $\mu_{\rm eff}$.

The chargino sector of these three models is the same except that in the NMSSM/nMSSM the parameter μ is replaced by $\mu_{\rm eff}$. The charginos $\tilde{\chi}_{1,2}^{\pm}$ ($m_{\chi_1^\pm} \leq m_{\chi_2^\pm}$) are the mixture of charged Higgsinos $\psi_{H_{u,d}}^{\pm}$ and winos $\lambda^{\pm} = (\lambda^1 \pm \lambda^2)/\sqrt{2}$, whose mass matrix in the basis of $(-i\lambda^{\pm}, \psi_{H_{u,d}}^{\pm})$ is given by

$$\begin{pmatrix} M_2 & \sqrt{2}m_W s_b \\ \sqrt{2}m_W c_b & \mu_{\rm eff} \end{pmatrix}. \tag{13}$$

So the chargino $\tilde{\chi}_1^{\pm}$ can be wino-dominant (when M_2 is much smaller than μ) or higgsino-dominant (when μ is much smaller than M_2). Since the composing property (wino-like, bino-like, higgsino-like or singlino-like) of the LSP and the chargino $\tilde{\chi}_1^{\pm}$ is very important in SUSY phenomenologies, we will show such a property in our following study.

III. COMPARISON WITH THE MSSM AND THE MSSM WITH A SINGLET SECTOR

A. In light of CDMS II and XENON

First let's see the MSSM, the NMSSM and the nMSSM under the constraints of results of CDMS II and XENON100. As both current and future limits of $\tilde{\chi}$-nucleon of CDMS and XENON are similar to each other, we will show only one of them. Nevertheless, as a good substitute of the SM, SUSY model must satisfy all the results of current collider and detector measurements. In our study we consider the following experimental constraints: [15] (1) we require $\tilde{\chi}_1^0$ to account for dark matter relic density $0.105 < \Omega h^2 < 0.119$; (2) we require the SUSY contribution to explain the deviation of the muon a_μ, i.e., $a_\mu^{\rm exp} - a_\mu^{\rm SM} = (25.5 \pm 8.0) \times 10^{-10}$, at 2σ level; (3) the LEP-I bound on the invisible Z-decay, $\Gamma(Z \to \tilde{\chi}_1^0\tilde{\chi}_1^0) < 1.76$

MeV, and the LEP-II upper bound on $\sigma(e^+e^- \to \tilde{\chi}^0_1\tilde{\chi}^0_i)$, which is 5×10^{-2} pb for $i > 1$, as well as the lower mass bounds on sparticles from direct searches at LEP and the Tevatron; (4) the constraints from the direct search for Higgs bosons at LEP-II, including the decay modes $h \to h_1h_1, a_1a_1 \to 4f$, which limit all possible channels for the production of the Higgs bosons; (5) the constraints from B physics observable such as $B \to X_s\gamma$, $B_s \to \mu^+\mu^-$, $B^+ \to \tau^+\nu$, $\Upsilon \to \gamma a_1$, the a_1–η_b mixing and the mass difference ΔM_d and ΔM_s; (6) the constraints from the precision EW observable such as $\rho_{\rm lept}$, $\sin^2\theta_{\rm eff}^{\rm lept}$, m_W and R_b; (7) the constraints from the decay $\Upsilon \to \gamma h_1$, and the Tevatron search for a light Higgs boson via 4μ and $2\mu2\tau$ signals [16]. The constraints (1–5) have been encoded in the package NMSSMTools [17]. We use this package in our calculation and extend it by adding the constraints (6, 7). As pointed out in Ref. [16], the constraints (7) are important for a light Higgs boson. In addition to the above experimental limits, we also consider the constraint from the stability of the Higgs potential, which requires that the physical vacuum of the Higgs potential with non-vanishing VEVs of Higgs scalars should be lower than any local minima.

For the calculation of cross section of $\tilde{\chi}$-nucleon scattering, we use the formulas in Ref. [18, 19] for the MSSM and extend them to the NMSSM/nMSSM. It is sufficient to consider only the SI interactions between $\tilde{\chi}^0_1$ and nucleon (denoted by f_p for proton and f_n for neutron [19]) in the calculation. The leading order of these interactions are induced by exchanging the SM-like Higgs boson at tree level. For moderately light Higgs bosons, f_p is approximated by [19] (similarly forf_n)

$$f_p \simeq \sum_{q=u,d,s} \frac{f_q^H}{m_q} m_p f_{T_q}^{(p)} + \frac{2}{27} f_{T_G} \sum_{q=c,b,t} \frac{f_q^H}{m_q} m_p \ , \tag{14}$$

where $f_{Tq}^{(p)}$ denotes the fraction of m_p (proton mass) from the light quark q while $f_{T_G} = 1 - \sum_{u,d,s} f_{T_q}^{(p)}$ is the heavy quark contribution through gluon exchange. f_q^H is the coefficient of the effective scalar operator. The $\tilde{\chi}^0$-nucleus scattering rate is then given by [19]

$$\sigma^{SI} = \frac{4}{\pi} \left(\frac{m_{\tilde{\chi}^0} m_T}{m_{\tilde{\chi}^0} + m_T} \right)^2 \times \left(n_p f_p + n_n f_n \right)^2, \tag{15}$$

where m_T is the mass of target nucleus and $n_p(n_n)$ is the number of proton (neutron) in the target nucleus. In our numerical calculations we take $f_{T_u}^{(n)} = 0.023$, $f_{T_d}^{(n)} = 0.034$, $f_{T_u}^{(p)} = 0.019$, $f_{T_d}^{(p)} = 0.041$ and $f_{T_s}^{(p)} = f_{T_s}^{(n)} = 0.38$. Note that the scattering rate is very sensitive to the value of f_{T_s} [21]. Recent lattice simulation [22] gave a much smaller value of f_{T_s} (0.020), it reduces the scattering rate significantly which can be seen in Ref. [23].

Considering all the constraints listed above, we scan over the parameters in the following ranges

$$
\begin{aligned}
&100 \text{ GeV} \leq (M_{\tilde{q}}, M_{\tilde{\ell}},\ m_A,\ \mu) \leq 1 \text{ TeV}, \\
&50 \text{ GeV} \leq M_1 \leq 1 \text{ TeV}, \quad 1 \leq \tan\beta \leq 40, \\
&(|\lambda|, |\kappa|) \leq 0.7, \quad |A_\kappa| \leq 1 \text{ TeV},
\end{aligned} \tag{16}
$$

where $M_{\tilde{q}}$ and $M_{\tilde{\ell}}$ are the universal soft mass parameters of the first two generations of squarks and the three generations of sleptons respectively. To reduce the number of the relevant soft parameters, we worked in the so-called m_h^{max} scenario with following choice of the soft masses for the third generation squarks: $M_{\tilde{Q}_3} = M_{\tilde{U}_3} = M_{\tilde{D}_3} = 800$ GeV, and $X_t = A_t - \mu \cot\beta = -1600$ GeV. The advantage of such a choice is that other SUSY parameters more easily survive the constraints (so that the bounds we obtain are conservative). Moreover, we assume the grand unification relation for the gaugino masses:

$$
M_1 : M_2 : M_3 \simeq 1 : 1.83 : 5.26\ . \tag{17}
$$

This relation is often assumed in studies of SUSY at the TeV scale for it can be easily generated in the mSUGRA model [24]. Note that relaxing this relation will give a large effect on the light neutralino scenario [25].

The surviving points for the three model are displayed in Fig. 1 for the spin-independent elastic cross section of $\tilde{\chi}$-nucleon scattering. We see that for each model the CDMS II limits can exclude a large part of the parameter space allowed by current collider constraints and the future SuperCDMS (25 kg) limits can cover most of the allowed parameter space. For the MSSM and the NMSSM dark matter mass is roughly in range of 50-400 GeV, while for the nMSSM dark matter mass is constrained below 40 GeV by current experiments and further constrained below 20GeV by SuperCDMS in case of non-observation.

From Fig. 1, we can see that the $\tilde{\chi}$-nucleon scattering plot of the MSSM and the NMSSM are very similar to each other, but very different from nMSSM. This implies that under the experiment constraints, the singlet sector will decouple from the MSSM-like sector in the NMSSM, then the NMSSM will perform almost the same as the MSSM, However, the singlet components change EW scale phenomenology greatly in the nMSSM. This can also be seen in Fig. 2 and Fig. 3. We can see that for both the MSSM and the NMSSM $\tilde{\chi}_1^0$ is bino-dominant, while for the nMSSM $\tilde{\chi}_1^0$ is singlino-dominant, and the region allowed by

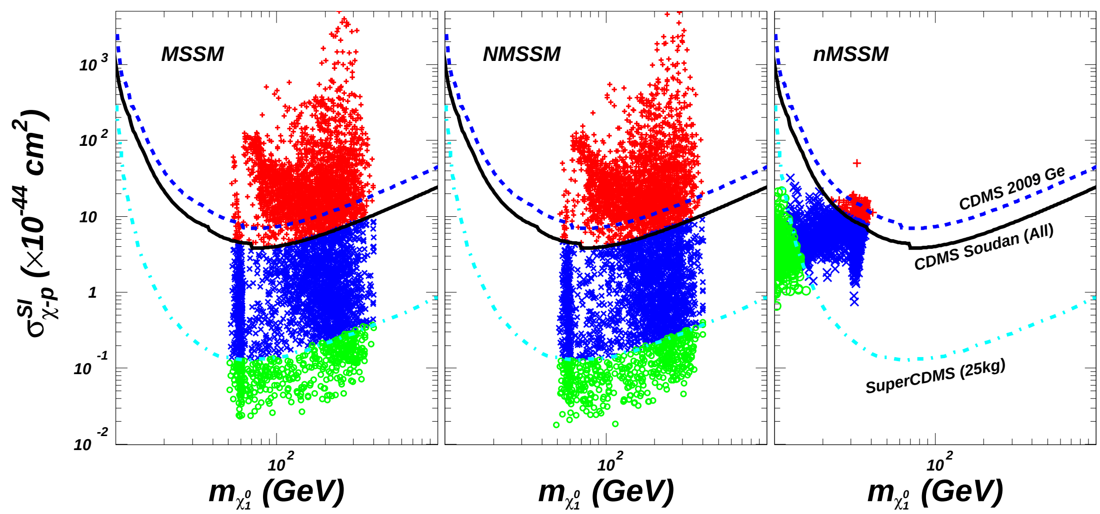

FIG. 1: The scatter plots (taken for Ref. [12]) for the spin-independent elastic cross section of $\tilde{\chi}$-nucleon scattering. The '+' points (red) are excluded by CDMS limits (solid line), the '×' (blue) would be further excluded by SuperCDMS 25kg [20] in case of non-observation (dash-dotted line), and the '○' (green) are beyond the SuperCDMS sensitivity.

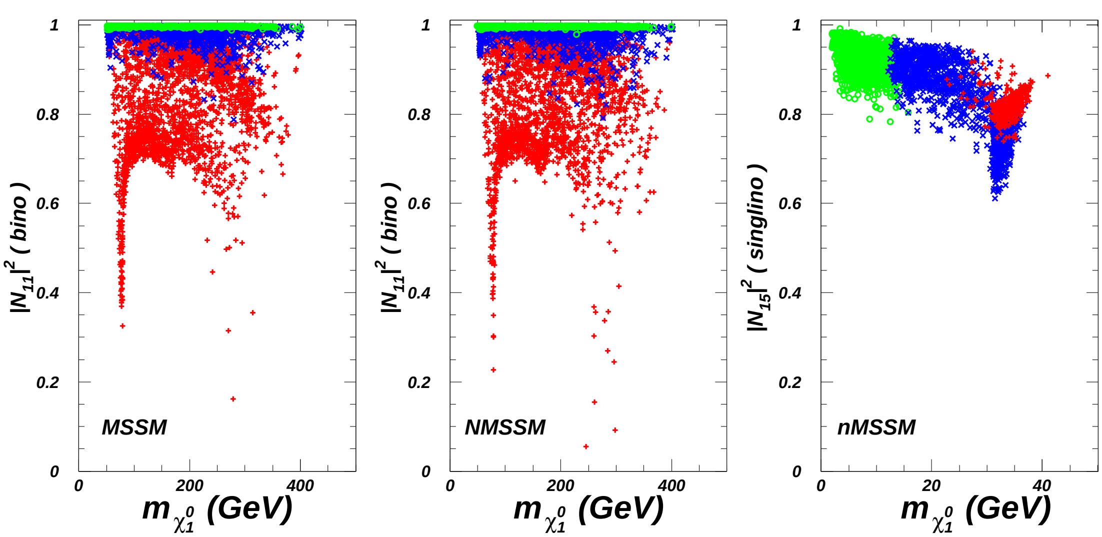

FIG. 2: Same as Fig. 1, but projected on the plane of $|N_{11}|^2$ and $|N_{15}|^2$ versus dark matter mass. (taken for Ref. [12])

CDMS limits (and SuperCDMS limits in case of non-observation) favors a more bino-like $\tilde{\chi}^0_1$ for the MSSM/NMSSM and a more singlino-like $\tilde{\chi}^0_1$ for the nMSSM. For the MSSM/NMSSM the LSP lower bound around 50 GeV is from the chargino lower bound of 103.5 GeV plus the assumed GUT relation $M_1 \simeq 0.5 M_2$; while the upper bound around 400 GeV is from the

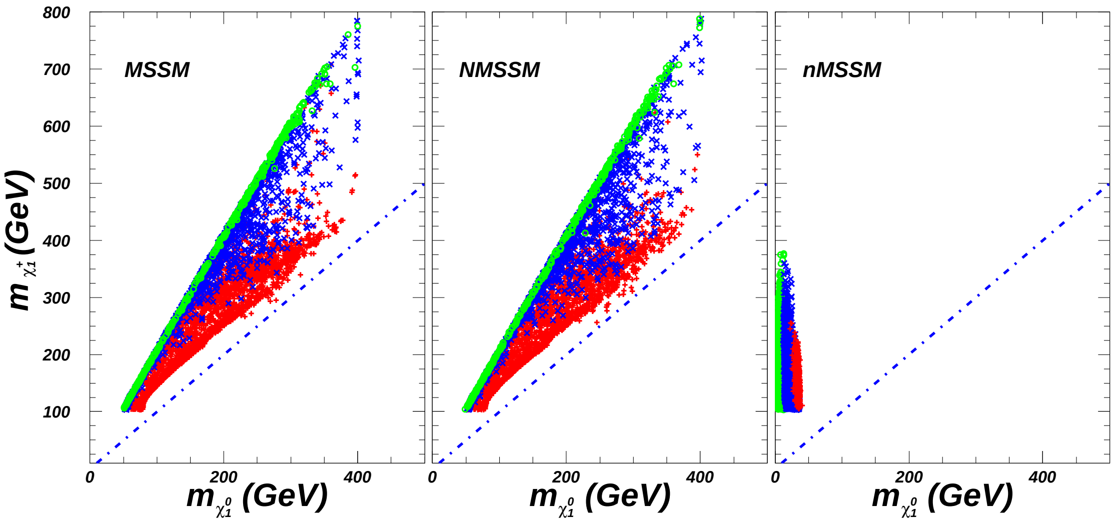

FIG. 3: Same as Fig. 1, but showing the chargino mass $m_{\chi_1^+}$ versus the LSP mass. The dashed lines indicate $m_{\chi_1^+} = m_{\chi_1^0}$. (taken for Ref. [12])

bino nature of the LSP (M_1 cannot be too large, must be much smaller than other relevant parameters) plus the experimental constraints like the muon g-2 and B-physics. If we do not assume the GUT relation $M_1 \simeq 0.5 M_2$, then M_1 can be as small as 40 GeV and the LSP lower bound in the MSSM/NMSSM will not be sharply at 50 GeV. (We talk about it in the following section.) For both the MSSM and the NMSSM, the CDMS limits tend to favor a heavier chargino and ultimately the SuperCDMS limits tend to favor a wino-dominant chargino with mass about $2m_{\chi_1^0}$. Note that, there still can be a singlino dominant LSP in some parameter space of the NMSSM [26], but in the scan range Eq. (16) listed above, getting such singlino dominant LSP needs some fine-tuning, thus we do not focus on it.

In Fig. 4 we show the value of $|\lambda|$ versus the charged Higgs mass in the NMSSM and the nMSSM. This figure indicates that λ larger than 0.4 is disfavored by the NMSSM. The underlying reason is that $h_1 \tilde{\chi}_1^0 \tilde{\chi}_1^0$ depends on λ explicitly and large λ can enhance $\tilde{\chi}$-nucleon scattering rate. By contrast, although CDMS has excluded some points with large λ in the nMSSM, there are still many surviving points with λ as large as 0.7. We have talked the reason above: to get a heavy χ_1^0, one need a large λ, a small $\tan\beta$ as well as a moderate μ_{eff}.

From the survived parameter space for all the model above, we should know that the Higgs decay will be similar for the MSSM and the NMSSM, but quite different from the nMSSM. This can be seen in Fig. 5 which shows decay branching ratio of $h_1 \to \tilde{\chi}_1^0 \tilde{\chi}_1^0$ versus the mass of the SM-like Higgs boson h_{SM} (which is h_1 here, and it is Higgs doublet $\hat{H}_u$ and

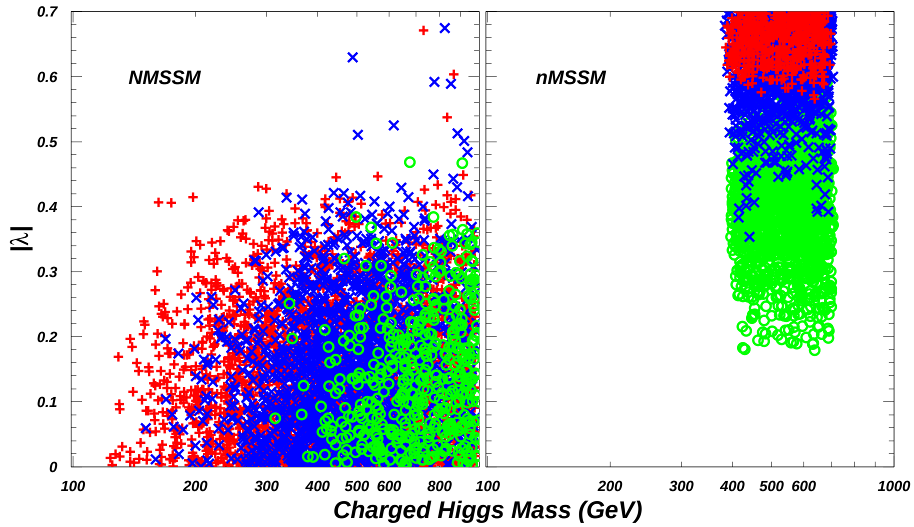

FIG. 4: Same as Fig. 1, but projected on the plane of $|\lambda|$ versus the charged Higgs mass in the NMSSM and the nMSSM. (taken for Ref. [12])

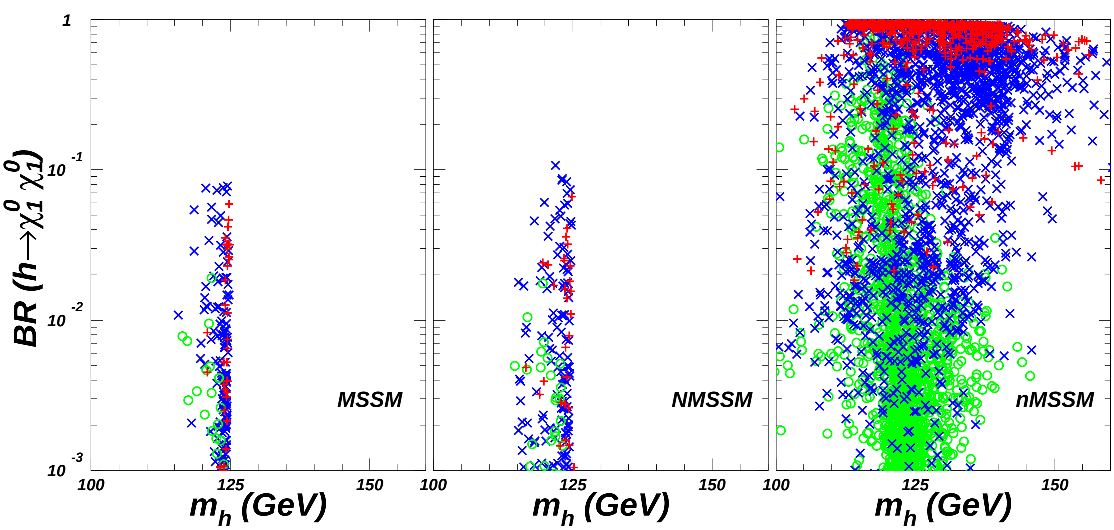

FIG. 5: Same as Fig. 1, but projected for the decay branching ratio of $h_{\rm SM} \to \chi_1^0\chi_1^0$ versus the mass of the Higgs boson $h_{\rm SM}$. (taken for Ref. [12])

$\hat{H}_d$ dominant). Such a decay is strongly correlated to the $\tilde{\chi}$-nucleon scattering because the coupling $h_1\tilde{\chi}_1^0\tilde{\chi}_1^0$ is involved in both processes. We see that in the MSSM and the NMSSM this decay mode can open only in a very narrow parameter space since $\tilde{\chi}_1^0$ cannot be so light, and in the allowed region this decay has a very small branching ratio (below 10%). However, in the nMSSM this decay can open in a large part of the parameter space since the LSP can be very light, and its branching ratio can be quite large (over 80% or 90%).

B. light dark matter in the NMSSM

As talked in the introduction, the data of CoGeNT experiment favors a light dark matter around 10 GeV. However, we scan the parameter space in the MSSM and find that it is very difficult to find a neutralino $\tilde{\chi}_1^0$ lighter than about 28 GeV, unless when it is associated with a light stau as the next to the lightest supersymmetric particle (NLSP), but such scenario always needs a fine-tuning in the parameter space [27]. The main reason for the absence of a lighter $\tilde{\chi}_1^0$ is that the dominant annihilation channel for $\tilde{\chi}_1^0$ in the early universe is $\tilde{\chi}_1^0\tilde{\chi}_1^0 \to b\bar{b}$ through s-channel exchange of the pseudoscalar Higgs boson (A) and the measured dark matter relic density requires $m_A \sim$ (90–100) GeV and $\tan\beta \sim 50$, this is in conflict with the constraints from the LEP experiment and B physics [28–30]. The LHC data gives an even more stronger constraint on the light pseudoscalar scenario [31] such that light dark matter seems impossible in the MSSM. Though in the nMSSM the neutralino $\tilde{\chi}_1^0$ can be as light as 10 GeV (shown in Fig. 1), the scattering rate is much lower under the CoGeNT-favored region. In the NMSSM, however, with the participation of singlet sector one can get very light [4] Higgs. This feature is particularly useful for light $\tilde{\chi}_1^0$ scenario since it opens up new important annihilation channels for $\tilde{\chi}_1^0$, i.e., either into a pair of h_1 (or a_1) or into a pair of fermions via s-channel exchange of h_1 (or a_1) [30, 32, 33]. For the former case, $\tilde{\chi}_1^0$ must be heavier than h_1 (a_1); while for the latter case, due to the very weak couplings of h_1 (a_1) with $\tilde{\chi}_1^0$ and with the SM fermions, a resonance enhancement (i.e. m_{h_1} or m_{a_1} must be close to $2m_{\tilde{\chi}_1^0}$) is needed to accelerate the annihilation. So a light $\tilde{\chi}_1^0$ may be necessarily accompanied by a light h_1 or a_1 to provide the required dark matter relic density. From the discussion in the upper section, light $\tilde{\chi}_1^0$ can be obtained by releasing the GUT relation Eq. (17), thus LSP in the NMSSM may explain the detection of CoGeNT. Note that, as the LSP in the nMSSM is singlino dominant, relaxing the GUT relation will not the change the phenomenology of dark matter and Higgs too much.

Now we discuss how to get a light h_1 or a_1 in the NMSSM. A light a_1 can be easily obtained when the theory is close to the $\mathrm{U}(1)_R$ or $\mathrm{U}(1)_{\mathrm{PQ}}$ symmetry limit, which can be realized by setting the product κA_κ to be negatively small [4]. In contrast, a light h_1 can not be obtained easily. However, as shown below, it can still be achieved by somewhat subtle cancelation via tuning the value of A_κ. We note that for any theory with multiple Higgs fields, the existence of a massless Higgs boson implies the vanishing of the determinant of

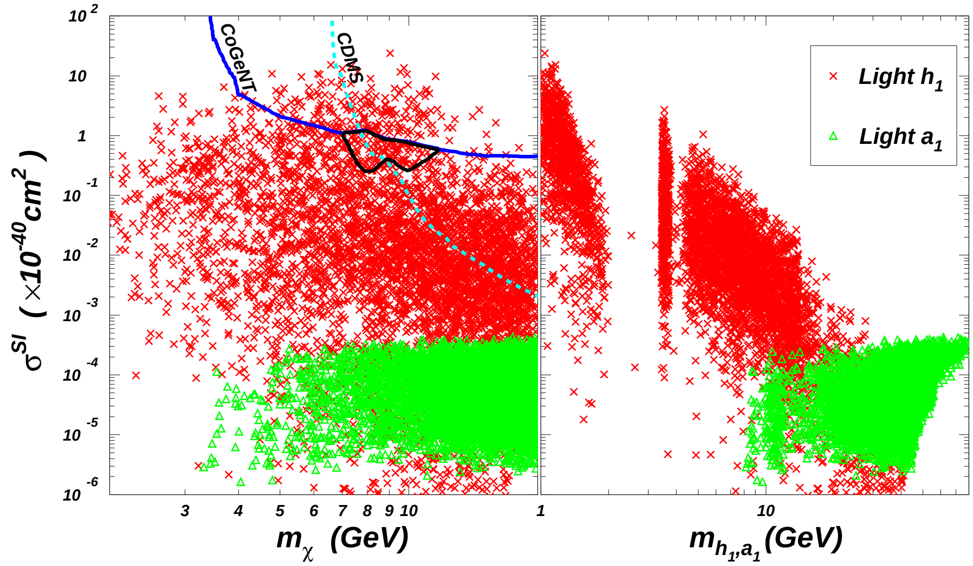

FIG. 6: The scatter plots (taken for Ref. [13]) of the parameter samples which survive all constraints, with '×' (red) and '▲' (green) corresponding to a light h_1 and a light a_1, respectively. The left frame is projected on the $\sigma^{\rm SI}$-m_χ plane, while the right frame is projected on the $\sigma^{\rm SI}$-m_{h_1} plane (denoted by '×') and the $\sigma^{\rm SI}$-m_{a_1} plane (denoted by '▲'). The curves are the limits from CoGeNT [9], CDMS [7], while the contour is the CoGeNT-favored region [9].

its squared mass matrix and vice versa. For the NMSSM, at tree level the parameter A_κ only enters the mass term of the singlet Higgs bosons, so the determinant ($\mathrm{Det}\,\mathcal{M}^2$) of the mass matrix of the CP-even Higgs bosons depends on A_κ linearly [4]. When other relevant parameters are fixed, one can then obtain a light h_1 by varying A_κ around the value $\tilde{A}_\kappa$ which is the solution to the equation $\mathrm{Det}\,\mathcal{M}^2 = 0$. In practice, one must include the important radiative corrections to the Higgs mass matrix, which will complicate the dependence of $\mathcal{M}^2$ on A_κ. However, we checked that the linear dependence is approximately maintained by choosing the other relevant parameters at the SUSY scale, and one can solve the equation iteratively to get the solution $\tilde{A}_\kappa$.

In Fig. 6 we display the surviving parameter samples, showing the $\tilde{\chi}$-nucleon scattering cross section versus the neutralino dark matter mass (left frame) and versus the mass of h_1 or a_1 (right frame). It shows that the scattering rate of the light dark matter can reach the sensitivity of CDMS and, consequently, a sizable parameter space is excluded by the CDMS data [20]. The future CDMS experiment can further explore (but cannot completely cover) the remained parameter space. Note that in the light-h_1 case the scattering rate

can be large enough to reach the sensitivity of CoGeNT and to cover the CoGeNT-favored region. The underlying reason is that the $\tilde{\chi}$-nucleon scattering can proceed through the t-channel exchange of the CP-even Higgs bosons, which can be enhanced by a factor $1/m_{h_1}^4$ for a light h_1 [32]; while a light a_1 can not give such an enhancement because the CP-odd Higgs bosons do not contribute to the scattering in this way. We noticed that the studies in [29, 34] claimed that the NMSSM is unable to explain the CoGeNT data because they did not consider the light-h_1 case.

In the light $\tilde{\chi}_1^0$ scenario, $h_{\rm SM}$ may decay exotically into $\tilde{\chi}_i^0\tilde{\chi}_j^0$, h_1h_1 or a_1a_1, and consequently the conventional decays are reduced. This feature is illustrated in Fig. 7, which shows that the sum of the exotic decay branching ratios may exceed 50% and the traditional decays $h_{\rm SM} \to b\bar{b}, \tau\bar{\tau}, WW^*, \gamma\gamma$ can be severely suppressed. Numerically, we find that the branching ratio of $h_{\rm SM} \to b\bar{b}$ is suppressed to be below 30% for all the surviving samples in the light-h_1 (h_2 is $h_{\rm SM}$) case and for about 96% of the surviving samples in the light-a_1 (h_1 is $h_{\rm SM}$) case (for the remaining 4% of the surviving samples in the light-a_1 case, the decay $h_{\rm SM} \to a_1a_1$ is usually kinematically forbidden so that the ratio of $h_{\rm SM} \to b\bar{b}$ may exceed 60%). Another interesting feature shown in Fig. 7 is that, due to the open-up of the exotic decays, $h_{\rm SM}$ may be significantly lighter than the LEP bound. This situation is favored by the fit of the precision electro-weak data and is of great theoretical interest [35].

Since the conventional decay modes of $h_{\rm SM}$ may be greatly suppressed, especially in the light-h_1 case which can give a rather large $\tilde{\chi}$-nucleon scattering rate, the LHC search for $h_{\rm SM}$ via the traditional channels may become difficult. Now the LHC observed a new particle in the mass region around 125-126 GeV which is the most probable the long sought Higgs boson [36]. In this mass range, the most important discovering channel of $h_{\rm SM}$ at the LHC is the di-photon signal. In Fig. 8 we give the ratio of the di-photon production rate to the SM at the LHC with $\sqrt{s} = 7$ TeV. In calculating the rate, we used the narrow width approximation and only considered the leading contributions to $pp \to h_{\rm SM}$ from top quark, bottom quark and the squark loops.

Fig. 8 indicates that, compared with the SM prediction, the ratio in the NMSSM in the light $\tilde{\chi}_1^0$ scenario is suppressed to be less than 0.4 for the light-h_1 case. For the light-a_1 case, most samples (about 96%) predict the same conclusion. Since in the light-h_1 case the $\tilde{\chi}$-nucleon scattering rate can reach the CoGeNT sensitivity, this means that in the framework of the NMSSM the CoGeNT search for the light dark matter will be correlated with the

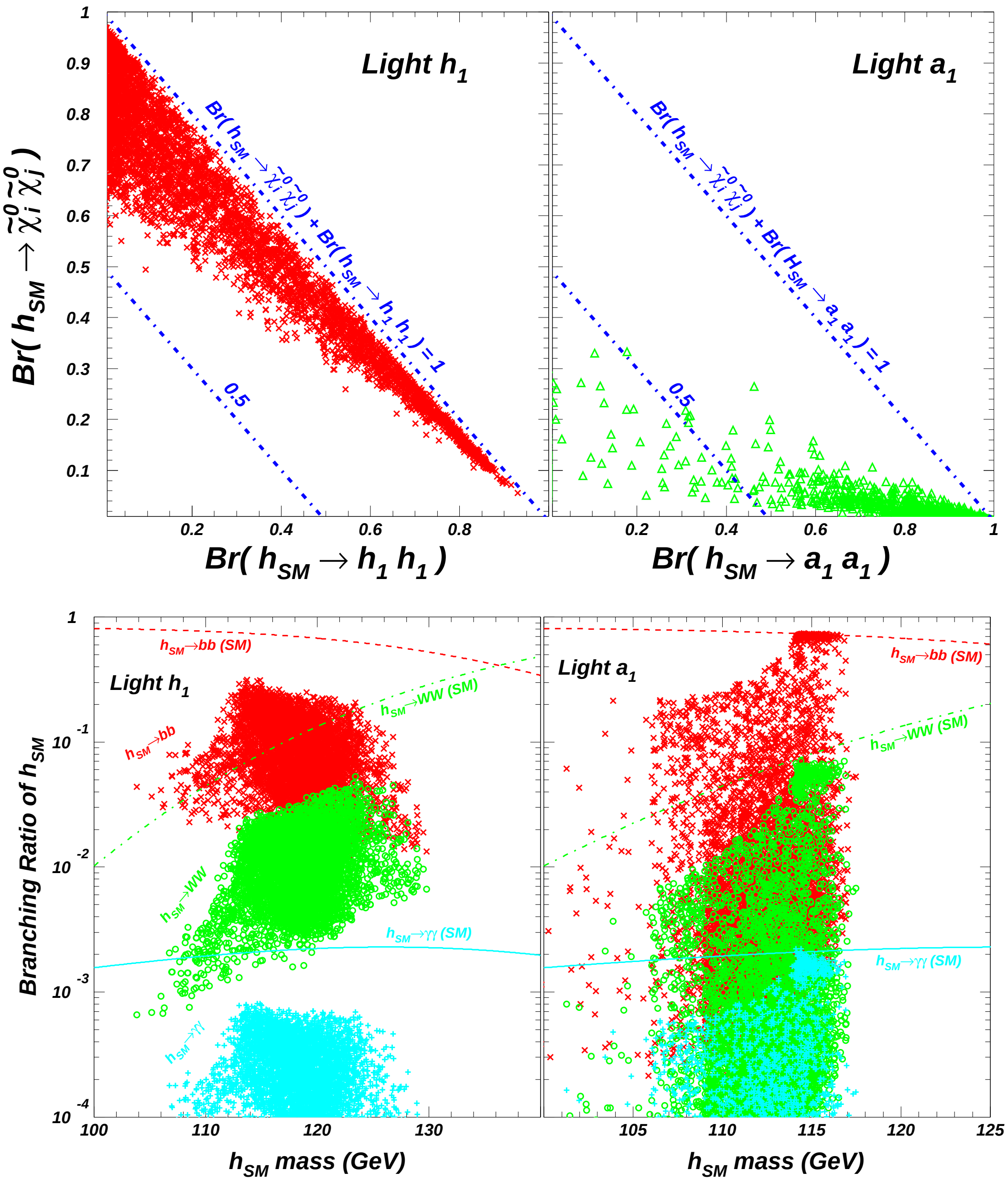

FIG. 7: Same as Fig. 6, but showing the decay branching ratios of the SM-like Higgs boson $h_{\rm SM}$. Here $Br(h_{\rm SM} \to \tilde{\chi}_i^0 \tilde{\chi}_j^0)$ denotes the total rates for all possible $h_{\rm SM} \to \tilde{\chi}_i^0 \tilde{\chi}_j^0$ decays. (taken for Ref. [13])

LHC search for the Higgs boson via the di-photon channel. We checked that, once the future XENON experiment fails in observing dark matter, less than 1% of the surviving samples in light a_1 case predict the ratio of di-photon signal larger than 0.4.

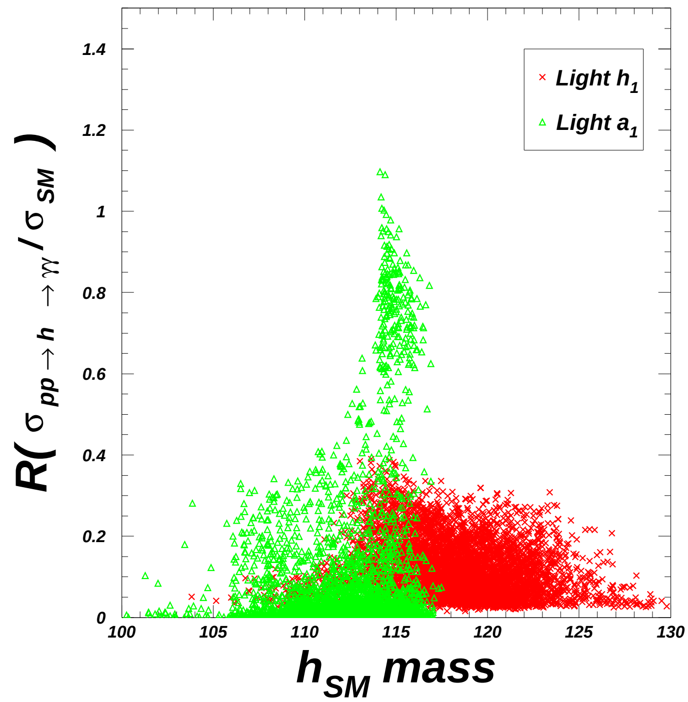

FIG. 8: Same as Fig. 6, but showing the diphoton production rate of the SM-like Higgs boson at the LHC.

IV. GENERAL EXTENSION FOR THE EXPLANATION TO PAMELA

To explain the PAMELA excess by dark matter annihilation, there are some challenges. First, dark matter must annihilate dominantly into leptons since PAMELA has observed no excess of anti-protons [11] (However, as pointed in Ref. [37], this statement may be not so solid due to the significant astrophysical uncertainties associated with their propagation). Second, the explanation of PAMELA excess requires an annihilation rate which is too large to explain the relic abundance if dark matter is produced thermally in the early universe. To tackle these difficulties, a new theory of dark matter was proposed in Ref. [38]. In this new theory the Sommerfeld effect of a new force in the dark sector can greatly enhance the annihilation rate when the velocity of dark matter is much smaller than the velocity at freeze-out in the early universe, and dark matter annihilates into light particles which are kinematically allowed to decay to muons or electrons.

The above fancy idea is hard to realize in the MSSM, because there is not a new force in the neutralino dark matter sector to induce the Sommerfeld enhancement and neutralino dark matter annihilates largely to final states consisting of heavy quarks or gauge and/or Higgs bosons [19, 39]. However, as discussed in Ref. [40], in a general extension of the

MSSM by introducing a singlet Higgs superfield, the idea in Ref. [38] can be realized by the singlino-like neutralino dark matter:

(i) The singlino dark matter annihilates to the light singlet Higgs bosons and the relic density can be naturally obtained from the interaction between singlino and singlet Higgs bosons.

(ii) The singlet Higgs bosons, not related to electro-weak symmetry breaking, can be light enough to be kinematically allowed to decay dominantly into muons or electrons through the tiny mixing with the Higgs doublets.

(iii) The Sommerfeld enhancement needed in dark matter annihilation for the explanation of PAMELA result can be induced by the light singlet Higgs boson.

In the following section, we will show how does this happen, the Higgs decay are also investigated.

A. Higgs and neutralinos spectrum

If introduce a singlet Higgs to the MSSM in general, the renormalizable holomorphic superpotential of Higgs is given by Ref. [40]

$$W = \mu \widehat{H}_u \cdot \widehat{H}_d + \lambda \widehat{S} \widehat{H}_u \cdot \widehat{H}_d + \eta \widehat{S} + \frac{1}{2} \mu_s \widehat{S}^2 + \frac{1}{3} \kappa \widehat{S}^3 \ , \tag{18}$$

which include linear term, quadratic term, cubic term of singlet superfield (like Wess-Zumino model [41]). Note that in such case, we do not require the singlet to solve the μ problem. The soft SUSY-breaking terms are given by

$$\begin{aligned} V_{\rm soft} \ =& \ \tilde{m}_u^2 |H_u|^2 + \tilde{m}_d^2 |H_d|^2 + \tilde{m}_s^2 |S|^2 \\ & + (B\mu H_u \cdot H_d + \lambda A_\lambda \ H_u \cdot H_d S + C\eta S + \frac{1}{2} B_s \mu_s S^2 + \frac{1}{3} \kappa A_\kappa \ S^3 + \text{h.c.}) \, . \end{aligned} \tag{19}$$

After the Higgs fields develop the VEVs v_u, v_d and s, i.e., we get the similar Higgs spectrum as the NMSSM and the nMSSM which is

(1) The CP-even Higgs mass matrix in the basis (ϕ_u, ϕ_d, σ) is given by

$$\mathcal{M}_{h,11} = g^2 v_u^2 + \cot\beta\left[\lambda s(A_\lambda + \kappa s + \mu_s) + B\mu\right], \tag{20}$$

$$\mathcal{M}_{h,22} = g^2 v_d^2 + \tan\beta\left[\lambda s(A_\lambda + \kappa s + \mu_s) + B\mu\right], \tag{21}$$

$$\mathcal{M}_{h,33} = \lambda(A_\lambda + \mu_s)\frac{v_u v_d}{s} - \lambda\frac{\mu}{s}(v_u^2 + v_d^2) + \kappa s(A_\kappa + 4\kappa s + 3\mu_s) - \frac{C\eta}{s}, \tag{22}$$

$$\mathcal{M}_{h,12} = (2\lambda^2 - g^2)v_u v_d - \lambda s(A_\lambda + \kappa s + \mu_s) - B\mu, \tag{23}$$

$$\mathcal{M}_{h,13} = 2\lambda(\mu + \lambda s)v_u - \lambda v_d(A_\lambda + 2\kappa s + \mu_s), \tag{24}$$

$$\mathcal{M}_{h,23} = 2\lambda(\mu + \lambda s)v_d - \lambda v_u(A_\lambda + 2\kappa s + \mu_s), \tag{25}$$

where $g^2 = (g_1^2 + g_2^2)/2$ with g_1 and g_2 being respectively the coupling constant of SU(2) and U(1) in the SM.

(2) The CP-odd Higgs mass matrix $\mathcal{M}_a$ is given by

$$\mathcal{M}_{a,11} = (\tan\beta + \cot\beta)[\lambda s(A_\lambda + \kappa s + \mu_s) + B\mu], \tag{26}$$

$$\begin{aligned}\mathcal{M}_{a,22} = {}& 4\lambda\kappa v_u v_d + \lambda(A_\lambda + \mu_s)\frac{v_u v_d}{s} - \lambda\frac{\mu}{s}(v_u^2 + v_d^2)\\ & -\kappa s(3A_\kappa + \mu_s) - \frac{C\eta}{s} - 2B_s\mu_s,\end{aligned} \tag{27}$$

$$\mathcal{M}_{a,12} = \lambda\sqrt{v_u^2 + v_d^2}\,(A_\lambda - 2\kappa s - \mu_s). \tag{28}$$

Note that here we have dropped the Goldstone mode, thus there left a 2×2 mass matrix in the basis $(\tilde{A}, \xi)$. and it can be diagonalized by an orthogonal 2×2 matrix P' and the physical CP-odd states a_i are given by (ordered as $m_{a_1} < m_{a_2}$)

$$a_1 = P'_{11}\tilde{A} + P'_{12}S_I = P'_{11}(\cos\beta\varphi_u + \sin\beta\varphi_d) + P'_{12}\xi, \tag{29}$$

$$a_2 = P'_{21}\tilde{A} + P'_{22}S_I = P'_{21}(\cos\beta\varphi_u + \sin\beta\varphi_d) + P'_{22}\xi. \tag{30}$$

(3) The charged Higgs mass matrix $\mathcal{M}_\pm$ in the basis $\left(H_u^+, H_d^+\right)$ is given by

$$\mathcal{M}_\pm = \left(\lambda s(A_\lambda + \kappa s + \mu_s) + B\mu + h_u h_d(\frac{g_2^2}{2} - \lambda^2)\right)\begin{pmatrix}\cot\beta & 1\\ 1 & \tan\beta\end{pmatrix}, \tag{31}$$

(4) The neutralino mass matrix is :

$$\mathcal{M}_0 = \begin{pmatrix} M_1 & 0 & m_Z s_W s_b & -m_Z s_W c_b & 0 \\ 0 & M_2 & -m_Z c_W s_b & m_Z c_W c_b & 0 \\ m_Z s_W s_b & -m_Z s_W s_b & 0 & -\mu & -\lambda v c_b \\ -m_Z s_W c_b & -m_Z c_W c_b & -\mu & 0 & -\lambda v s_b \\ 0 & 0 & -\lambda v c_b & -\lambda v s_b & 2\kappa s + \mu_s \end{pmatrix}. \tag{32}$$

B. Explanation of PAMELA and implication on Higgs decays

To explain the observation of PAMELA, a_1 is singlet-dominant, while h_1 is singlet-dominant and the next-to-lightest h_2 is doublet-dominant ($h_{\rm SM}$). We use the notation:

$$a \equiv a_1, \qquad h \equiv h_1, \qquad h_{\rm SM} \equiv h_2. \tag{33}$$

As discussed in Ref. [40], when the lightest neutralino $\tilde{\chi}_1^0$ in Eq. (11) is singlino-dominant, it can be a perfect candidate for dark matter. As shown in Fig. 9, such singlino dark matter annihilates to a pair of light singlet Higgs bosons followed by the decay $h \to aa$ (h has very small mixing with the Higgs doublets and thus has very small couplings to the SM fermions). In order to decay dominantly into muons, a must be light enough. Further, in order to induce the Sommerfeld enhancement, h must also be light enough. From the superpotential term $\kappa\hat{S}^3$ we know that the couplings $h\tilde{\chi}_1^0\tilde{\chi}_1^0$ and $a\tilde{\chi}_1^0\tilde{\chi}_1^0$ are proportional to κ. To obtain the relic density of dark matter, κ should be $\mathcal{O}(1)$. h, a are singlet-dominant and $\tilde{\chi}_1^0$ is singlino-dominant, this implies small mixing between singlet and doublet Higgs fields. From the superpotential in Eq.(18) we see that this means the mixing parameter λ must be small enough. On the other hand, from Eq. (22) and Eq. (27) lightness of h_1 and a_1 also require λ and other term approaching to zero. Therefore, in our scan we require parameters A_κ and B_s has the relation:

$$A_\kappa \sim \left(-4\kappa s - 3\mu_s + \frac{C\eta}{\kappa s^2}\right), \tag{34}$$

$$2B_s\mu_s \sim \left(-3A_\kappa \kappa s - \mu_s \kappa s - \frac{C\eta}{s}\right), \tag{35}$$

to realize light h_1 and a_1.

The numerical results of this model are displayed in different planes in Figs.10-12. We see from Fig. 10 that in the range $2m_\mu < m_a < 2m_\pi$, a decays dominantly into muons. It

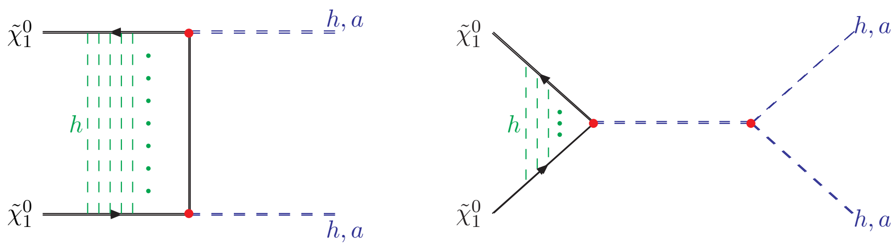

FIG. 9: Feynman diagrams for singlino dark matter annihilation where Sommerfeld enhancement is induced by exchanging h. (taken from Ref. [14])

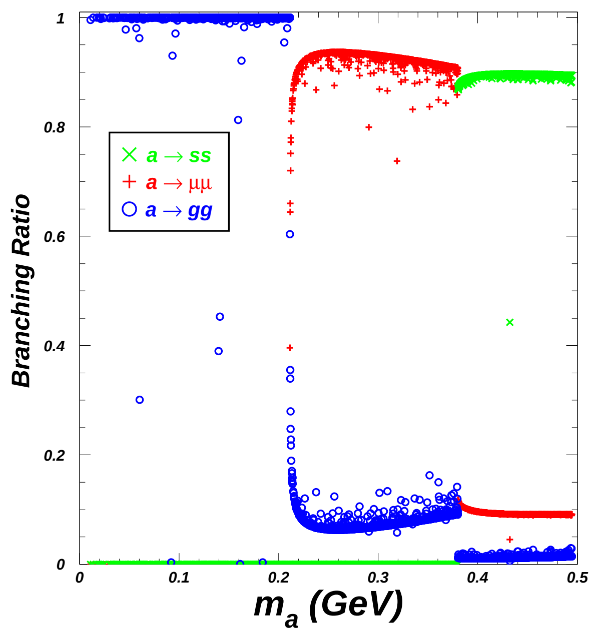

FIG. 10: The scatter plots showing the decay branching ratios $a \to \mu^+\mu^-$ (muon), $a \to gg$ (gluon) and $a \to s\bar{s}$ (s-quark) versus m_a for $\lambda = 10^{-3}$. (taken from Ref. [14])

is clear that h can be as light as a few GeV, which is light enough to induce the necessary Sommerfeld enhancement as shown in Fig. 11.

In left plot of Fig. 12, we show the branching ratios of $h_{\rm SM}$ decays. We see that in the allowed parameter space $h_{\rm SM}$ tends to decay into aa or hh instead of $b\bar{b}$. This can be understood as following, the MSSM parameter space is stringently constrained by the LEP experiments if $h_{\rm SM}$ is relatively light and decays dominantly to $b\bar{b}$, and to escape such stringent constraints $h_{\rm SM}$ tends to have exotic decays into aa or hh. As a result, the allowed parameter space tends to favor a large A_κ, as shown in right plot of Fig. 12, which greatly enhances the couplings $h_{\rm SM}aa$ and $h_{\rm SM}hh$ through the soft term $\kappa A_\kappa S^3$ although S has a small mixing with the doublet Higgs bosons. Such an enhancement can be easily seen. Take the coupling $h_{\rm SM}hh$ as an example, the soft term $\kappa A_\kappa S^3$ gives a term $\kappa A_\kappa \sigma^3$ which then

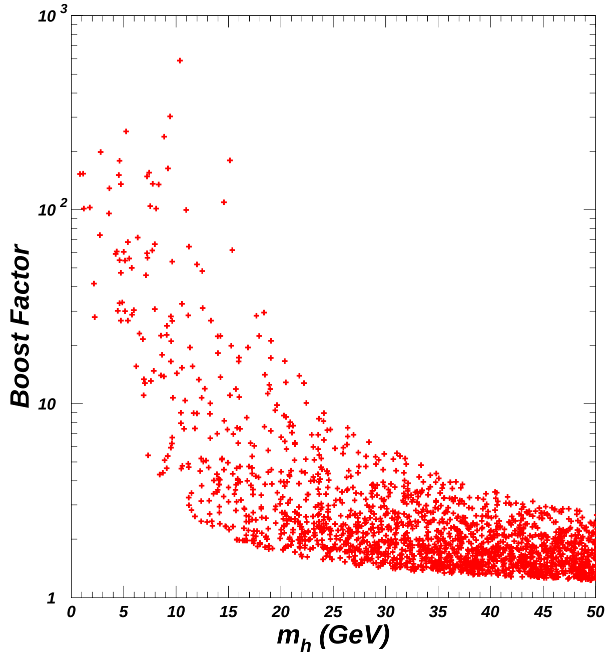

FIG. 11: Same as Fig. 10, but showing the Sommerfeld enhancement factor induced by h. (taken from Ref. [14])

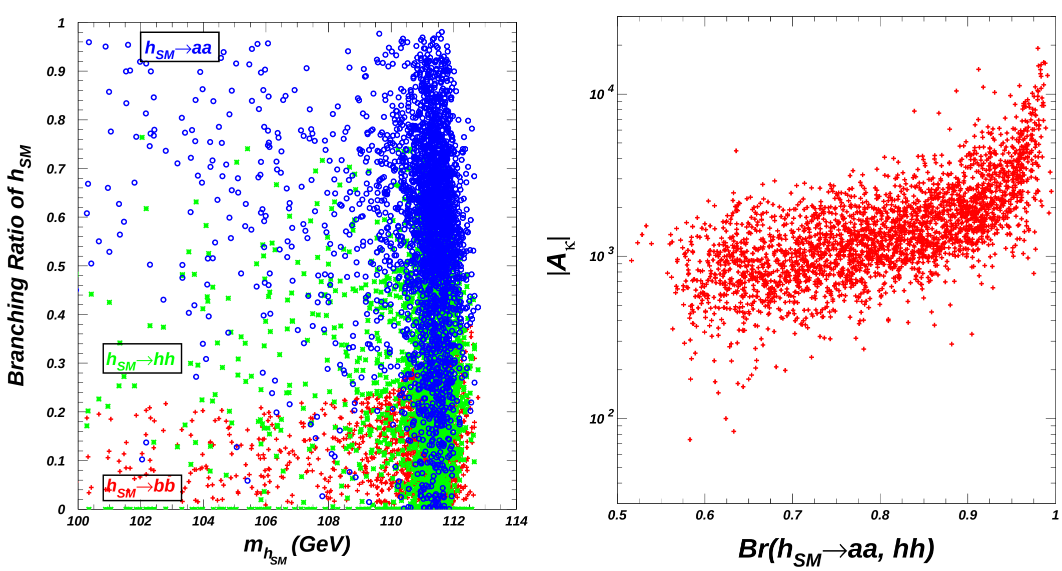

FIG. 12: Same as Fig. 10, but showing branching ratio of $h_{\rm SM} \to aa, hh$ versus $m_{h_{\rm sm}}$ and $|A_\kappa|$ versus the branching ratio of $h_{\rm SM} \to aa, hh$. (taken from Ref. [14])

gives the interaction $\kappa A_\kappa \ (U^H_{13})^2 U^H_{23} \ h_{\rm SM} hh$ because $\sigma = U^H_{13} h_1 + U^H_{23} h_2 + U^H_{33} h_3$ with $h_1 \equiv h$ and $h_2 \equiv h_{\rm SM}$ (see Eqs. (9) and (33)). Although the mixing $(U^H_{13})^2 U^H_{23}$ is small for a small λ, a large A_κ can enhance the coupling $h_{\rm SM} hh$. Note that as the mass of the observed Higgs boson at the LHC is around 125 GeV, thus in the MSSM, the dominant decay mode of $h_{\rm SM}$ is $b\bar{b}$. In this general singlet extension of the MSSM, its dominant decay mode may be

changed to aa or hh, as shown in our above results.

Finally, we note that for the specified singlet extensions like the nMSSM and the NMSSM, the explanation of PAMELA and relic density through Sommerfeld enhancement is not possible. The reason is that the parameter space of such models is stringently constrained by various experiments and dark matter relic density as shown in the above section, and, as a result, the neutralino dark matter may explain either the relic density or PAMELA, but impossible to explain both via Sommerfeld enhancement. For example, in the nMSSM various experiments and dark matter relic density constrain the neutralino dark matter particle in a narrow mass range [42], which is too light to explain PAMELA.

V. SUMMARY

At last we summarize here, the SUSY dark matter and Higgs physics will be changed if introducing a singlet to the MSSM. Under the latest results of dark matter detection, we have:

1. In the MSSM, the NMSSM and the nMSSM, the latest detection result can exclude a large part of the parameter space allowed by current collider constraints and the future SuperCDMS and XENON can cover most of the allowed parameter space.

2. Under the new dark matter constraints, the singlet sector will decouple from the MSSM-like sector in the NMSSM, thus the phenomenologies of dark matter and Higgs are similar to the MSSM. The singlet sector make the nMSSM quite different from the MSSM, the LSP in the nMSSM are singlet dominant, and the SM-like Higgs will mainly decay into the singlet sector. Future precision measurements will give us an opportunity to determine whether the new scalar is from standard model or from SUSY. Perhaps the nMSSM will be the first model be excluded for its much larger branching ratio of invisible Higgs decay.

3. The NMSSM can allow light dark matter at several GeV exists. Light CP-even or CP-odd Higgs boson must be present so as to satisfy the measured dark matter relic density. In case of the presence of a light CP-even Higgs boson, the light neutralino dark matter can explain the CoGeNT and DAMA/LIBRA results. Further, we find that in such a scenario the SM-like Higgs boson will decay predominantly into a pair

of light Higgs bosons or a pair of neutralinos and the conventional decay modes will be greatly suppressed.

4. The general singlet extension of the MSSM gives a perfect explanation for both the relic density and the PAMELA result through the Sommerfeld enhanced annihilation into singlet Higgs bosons (a or h followed by $h \to aa$) with a being light enough to decay dominantly to muons or electrons. Although the light singlet Higgs bosons have small mixing with the Higgs doublets in the allowed parameter space, their couplings with the SM-like Higgs boson h_{SM} can be enhanced by the soft parameter A_κ. In order to meet the stringent LEP constraints, the h_{SM} tends to decay into the singlet Higgs pairs aa or hh instead of $b\bar{b}$.

Acknowledgment

This work was supported in part by the NSFC No. 11005006, No. 11172008 and Doctor Foundation of BJUT No. X0006015201102.

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